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Aspects of Standard Models with Two Higgs Doublets [1]

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ABSTRACT

We [2] present some properties of the SM with two Higgs doublets. Unlike the minimal SM, the Yukawa couplings (in the usual "type II" model) converge with the increasing energy to flavor democracy (FD), i.e. to common values, in a specific flavor basis. This may represent a possible signal of some new, (almost) flavor-blind physics beyond the SM. When imposing the assumption of equality of the corresponding quark and leptonic Yukawa couplings at high transition energies, we can estimate the physical mass of the tau-neutrino as a function of m_t and the VEV ratio. Furthermore, such an assumption would effectively rule out the existence of the 4th generation of fermions.

We [3] also investigated the most general framework of the SM with two Higgs doublets such that no flavor-changing neutral currents (FCNC) occur at the tree level. Finite 1-loop-induced FCNC (and CP-violating) effects, when confronted with experimental constraints from the physics of K and B mesons, provide us with constraints on the values of the dominant Yukawa couplings of the charged Higgs with the top quark. In the usual, more restrictive, "type II" model, this would imply certain constraints on the value of the VEV ratio.

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1.) Flavor Democracy

We discuss in this part the behavior of different frameworks of the Standard Model with two Higgs doublets (2HD-SM) at high energies, with a view to the notion of the so called flavor democracy (FD). Two types of such models are frequently being used:

a) "type I" [2HD-SM(I)] - In this model just one Higgs doublet (say, $H^{(1)}$) couples to all fermions

$$\mathcal{L}^{(I)}_{Yukawa} = - \sum_{i,j=1}^{3} \{ D^{(q)}_{ij} (\bar{q}^{(i)}_L H^{(1)}) q^{(j)}_{dR} + U^{(q)}_{ij} (\bar{q}^{(i)}_L \tilde{H}^{(1)}) q^{(j)}_{uR} + \text{h.c.} \} + \cdots \ , \tag{1}$$

where the dots represent the analogous terms for the leptons. The following notations are used:

$$H = \begin{pmatrix} H^+ \\ H^0 \end{pmatrix} \ , \ \tilde{H} = i\tau_2 H^* \ , \quad q^{(i)} = \begin{pmatrix} q^{(i)}_u \\ q^{(i)}_d \end{pmatrix} \ , \quad q^{(1)} = \begin{pmatrix} u \\ d \end{pmatrix} \ , \ q^{(2)} = \begin{pmatrix} c \\ s \end{pmatrix} \ , \ q^{(3)} = \begin{pmatrix} t \\ b \end{pmatrix} \ ,$$

and similarly for the leptonic doublets $\ell^{(i)}$ containing Dirac neutrinos and charged leptons.

This model is very closely related to the minimal SM (MSM), the only difference in the Yukawa sector being the smaller vacuum expectation value $\langle (H^0)^{(1)} \rangle_o = v_1/\sqrt{2} < v/\sqrt{2}$ ($v \approx 246.22 GeV$), and hence the correspondingly larger Yukawa coupling parameters.

b) "type II" [2HD-SM(II)] - Here, one doublet ($H^{(1)}$) couples to the "down-type" right-handed fermions f_{dR} (q_{dR}, ℓ_{dR}) and is responsible for the "down-type" masses ($H^{(1)} \mapsto H^{(d)}$); the other doublet ($H^{(2)} \mapsto H^{(u)}$) couples to the "up-type" fermions f_{uR} (q_{uR}, ℓ_{uR}), being responsible for their masses:

$$\mathcal{L}^{(II)}_{Yukawa} = - \sum_{i,j=1}^{3} \{ D^{(q)}_{ij} (\bar{q}^{(i)}_L H^{(d)}) q^{(j)}_{dR} + U^{(q)}_{ij} (\bar{q}^{(i)}_L \tilde{H}^{(u)}) q^{(j)}_{uR} + \text{h.c.} \} + \cdots \ . \tag{2}$$

The mass matrices are proportional to the vacuum expectation values (VEVs) of the Higgses: $M^{(q,\ell)}_u = v_u U^{(q,\ell)}/\sqrt{2}$, $M^{(q,\ell)}_d = v_d D^{(q,\ell)}/\sqrt{2}$, where

$$\langle H^{(u)} \rangle_o = \frac{e^{i\xi}}{\sqrt{2}} \begin{pmatrix} 0 \\ v_u \end{pmatrix} \ , \qquad \langle H^{(d)} \rangle_o = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_d \end{pmatrix} \ , \qquad v_u^2 + v_d^2 = v^2 (\approx 246^2 GeV^2) \ .$$

The expressions are written in any flavor basis, i.e. a basis in which $f^{(i)}_L$ are $SU(2)_L$-isodoublets.

We note that the 2HD-SM(II), in comparison with the MSM and the closely related 2HD-SM(I), has essentially *different* (1-loop) renormalization group equations (RGEs) for the 3×3 Yukawa matrices

$$16\pi^2 \frac{dU^{(q)}}{d \ln E} = +\alpha D^{(q)} D^{(q)\dagger} U^{(q)} + \cdots \ , \qquad \text{etc.} \tag{3}$$

where $\alpha = 0.5$ in the 2HD-SM(II) ($\alpha = -1.5$ in the 2HD-SM(I) and MSM). E is here the running energy of probes. It turns out that this difference in the structure of the RGEs leads to a drastically different behavior of the corresponding models at high energies vis-à-vis the so called flavor democracy (FD), as will be shown in the following paragraphs.

We define that a model has the trend to FD (as the energy E increases toward the Landau pole Λ_{pole}) iff there exists a flavor basis (i.e. a basis in which f_L behave as $SU(2)_L$-isodoublets) such that the Yukawa couplings in the corresponding "up-type" and "down type" sectors of quarks (and leptons) converge to common values and the CKM-mixings converge to zero

$$U^{(q,\ell)} \to g^u_{q,\ell} A_{FD} \ , \qquad D^{(q,\ell)} \to g^d_{q,\ell} A_{FD} \ , \qquad V_{ckm} \to \mathbf{1} \qquad (\text{as } E \uparrow \Lambda_{pole}) \ . \tag{4}$$

Here, $\mathbf{1}$ is the 3×3 identity matrix and A_{FD} is the 3×3 flavor-democratic matrix

$$A_{FD} = \begin{bmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{bmatrix} .$$

The motivation behind this notion of FD-trend is the following. We could interpret the Landau pole Λ_{pole} (i.e. the energy where the Yukawa couplings diverge) as roughly the energy where the SM is replaced by a new physics ($\Lambda_{pole} \sim E_{trans.}$). Then the trend to FD, as defined above, would represent a signal that the physics beyond the SM is a strongly interacting physics with almost flavor-blind forces (where the tiny deviation from FD is provided by some even higher, unknown physics). However, the entire notion of the trend toward FD at high energies could also be regarded as independent of any specific motivation. On the other hand, the notion of FD at *low* energies is well known [1].

The FD matrix A_{FD} can be written in the diagonal ("mass") basis as the matrix with the diagonal elements $(0, 0, 3)$. Hence, when safely neglecting the very light 1st generation masses, we can rewrite the trend toward FD for the quark sectors (q-q FD) as

$$\frac{m_s}{m_b} \ , \ \frac{m_c}{m_t} \ , \ V_{cb} \ \to \ 0 \qquad (\text{as } E \uparrow \Lambda_{pole}) \ , \tag{5}$$

and analogously in the leptonic sectors (ℓ-ℓ FD), by replacing the quark masses by the corresponding masses of charged leptons and Dirac neutrinos. We note that already at low energies ($E \sim E_{ew}$), the SM is not far away from FD. By simply evolving the 1-loop RGEs for Yukawa couplings (no Higgs couplings involved) from low energies, where the "boundary conditions" are more or less known by experiments, to high energies, we arrive at the following results:

$$\text{MSM and 2HD-SM(I)} \ \not\to \ \text{FD} \ ; \qquad \text{2HD-SM(II)} \ \to \ \text{FD} \qquad (\text{as } E \uparrow \Lambda_{pole}) \ ,$$

as seen from Figs. 1 and 2.

Figs. 1, 2: Quark flavor democracy parameters m_s/m_b, m_c/m_t and V_{cb} as functions of the energy scale μ ($= E$), for the case of the minimal SM (left) and the 2HD-SM(II) with the VEV ratio $v_u/v_d = 0.5$ (right), for various top quark masses $m_t(m_t)$ (masses of leptons were ignored).

Figs. 1 and 2 are for the quark sectors (q-q FD), for various values of $m_t(E = m_t)$ ($\approx .96 m_t^{phy}$). Fig. 2 is for a specific ratio of vacuum expectation values (VEVs): $\tan\beta = v_u/v_d = 0.5$. In the MSM (and the 2HD-SM(I)), it is the "down-type" sector and the mixing that lead us away from FD as the energy increases (cf. Fig. 1). However, all the conclusions are the same also for the leptonic sectors (ℓ-ℓ FD), and for all other values of the VEVs. The convergence toward FD (in the 2HD-SM(II)) is the more striking (Λ_{pole} lower), the lower the VEV-ratio $\tan\beta$ and the higher m_t^{phy}. Specifically, for $m_t^{phy} = 175GeV$ and $\tan\beta \leq 0.5$, we have $\Lambda_{pole}(\sim E_{trans.}) \leq 33TeV$. The related fact is that choosing $m_t^{phy} \leq 200GeV$, we obtain $\Lambda_{pole}(\sim E_{trans.}) \lesssim E_{Planck}$ for all $\tan\beta \lesssim 1.75$, thus having for a large range of pertubatively allowed $\tan\beta$ a reasonable framework for condensation mechanism ($H^{(u)} \sim \langle t\bar{t}\cdots\rangle$, $H^{(d)} \sim \langle b\bar{b}+\cdots\rangle$), *unlike* the MSM and the closely related 2HD-SM(I) (cf. [2], where the condition $E_{condensation} \sim \Lambda_{pole} \lesssim E_{Planck}$ yields $m_t^{phy} \gtrsim 215GeV$). Condensation mechanisms leading from a Higgless higher physics to 2HD-SMs have been seriously investigated [3].

Up to this point, the results presented here can be regarded as formal properties of SMs with various Higgs sectors, formally independent of any assumptions about the physics beyong the SMs. However, in the FD-favored 2HD-SM(II), in order to reduce further the number of high energy Yukawa parameters, we can *choose* to impose the condition of the trend to the so called mixed quark-lepton FD (q-ℓ FD)

$$g_q^u \approx g_\ell^u \ , \qquad g_q^d \approx g_\ell^d \qquad \text{i.e.} \ \frac{m_\tau}{m_b} \ , \ \frac{m_{\nu_\tau^D}}{m_t} \ \approx \ 1 \qquad (\text{as } E \uparrow \Lambda_{pole}) \ . \tag{6}$$

Demanding this, and employing the (1-loop) RGEs, we obtain, for a given m_t^{phy} and $\tan\beta$, the values of the Dirac mass $m^o_{\nu_\tau^D}$ (at $E = 1GeV$) and Λ_{pole}. If employing in addition the usual see-saw mechanism [4], this would give us a good estimate (upper bound) for the physical neutrino mass

$$M_{Majorana} \gtrsim \Lambda_{pole} \ \Rightarrow$$

$$m_{\nu_\tau}^{phy} \simeq \frac{(m^o_{\nu_\tau^D})^2}{M_{Maj.}} \ \lesssim \frac{(m^o_{\nu_\tau^D})^2}{\Lambda_{pole}} \ (= (m_{\nu_\tau}^{phy})^{u.b.}) \ .$$

Demanding that this quantity not exceed the experimentally suggested upper value of $31MeV$, we obtain an upper bound for m_t^{phy} as a function of $\tan\beta$; similarly, the condition $\Lambda_{pole} < E_{Planck}$ provides us with the lower bound on m_t^{phy} (Fig. 3). For $m_t^{phy} = (175 \pm 20)GeV$, we obtain from Fig. 3: $0.53 < \tan\beta < 2.1$. These results, however, are obtained under the assumption of the imposed q-ℓ FD at $E \simeq \Lambda_{pole}$ (in the FD-favored 2HD-SM(II)).

Most of the results presented here up to this point have been published in refs.[5].

When imposing the q-ℓ FD (at $E \simeq \Lambda_{pole}$) in the 2HD-SM(II), we can ask ourselves whether such a framework would allow for the existence of the (heavy) 4th generation of fermions (t', b'), $(\nu^D_{\tau'}, \tau')$. In this case, the q-ℓ FD would mean: $m_{t'} \approx m_{\nu^D_{\tau'}}$, $m_{b'} \approx m_{\tau'}$ at $E \approx \Lambda_{pole}$. Similarly as before, for any chosen $m_{t'}^{phy}$ and $m_{b'}^{phy}$ (and for given 3rd generation masses m_t^{phy}, m_b^{phy} and m_τ^{phy}), we obtain by employing the 1-loop RGEs the correspondence

$$m_{t'}^{phy} \ , \ m_{b'}^{phy} \ \mapsto \ \Lambda_{pole} \ , \ (m_{\nu_{\tau'}}^{phy})^{u.b.} \ , \ m_{\tau'}^{phy} \ .$$

The correspondence is virtually independent of the Yukawa coupling of the 3rd generation neutrino ν_τ^D (at $E = 1GeV$). We applied see-saw, as before, to obtain the upper bound estimate $(m_{\nu_{\tau'}}^{phy})^{u.b.}$. However, at least 4 additional experimentally suggested constraints have to be imposed:

$$1.)\ m_{t'}^{phy}, m_{b'}^{phy} > m_t^{phy}\ ; \qquad 2.)\ 40GeV \leq m_{\nu_{\tau'}}^{phy}\ (\lesssim (m_{\nu_{\tau'}^D}^o)^2/\Lambda_{pole})\ ;$$

$$3.)\ \text{ the SM is perturbative (say: } m_{t'}^{phy}, m_{b'}^{phy} \lesssim 0.5\Lambda_{pole})\ ; \quad 4.)\ (\triangle\rho)_{\text{from heavy fermions}} \leq 0.0076.$$

The fourth constraint is from an essentially model-independent analysis of the LEP data [6]. It turns out that the four constraints above can be satisfied only if $m_t^{phy} < 155GeV$. Hence, in the light of the recent CDF measurement of m_t^{phy} ($\simeq (174 \pm 20)GeV$), we conclude that the existence of the 4th generation is practically ruled out within this framework (i.e. in the FD-favored 2HD-SM(II), with see-saw, and q-ℓ FD at Λ_{pole}) [7].

2.) Suppression of the flavor-changing neutral currents in the 2HD-SM(II)

Experiments show that the flavor-changing neutral currents (FCNC) are very suppressed

$$\frac{\Gamma(K_L^0 \to \mu^+\mu^-)}{\Gamma(K_L^0 \to all)} \simeq (7.3 \pm 0.4)10^{-9}\ , \qquad |m_{B_d^0} - m_{\bar{B}_d^0}| \simeq (3.6 \pm 0.7)10^{-10}MeV\ ,$$

$$|m_{K_L} - m_{K_S}| \simeq (3.52 \pm 0.02)10^{-12}MeV\ ,\ \text{etc.}$$

Therefore, the usual conditions imposed in various SMs are: a) *no* FCNC at the tree level, and b) the 1-loop-induced FCNC are sufficiently suppressed.

The conditions for the 1-loop FCNC suppression for gauge boson loops, i.e. the allowed representation contents of fermions, have been investigated some time ago [8]. In addition, Glashow and Weinberg [8] proposed for the Higgs sector the MSM (i.e. one Higgs doublet) and "type I" and/or "type II" 2HD-SM (two Higgs doublets; cf. eqs. (1) and (2)) - as models which, in a "natural" way, have no FCNC in the Yukawa couplings at the tree level.

However, recently Y.-L. Wu [9] has pointed out the most general framework of the 2HD-SMs not having FCNC at the tree level in the Yukawa sector ("type III"). The Yukawa interactions in this framework, in any flavor basis, have the most general form

$$\begin{aligned} \mathcal{L}_{Y.}^{(III)} &= -\sum_{i,j=1}^{3} \{ D_{ij}^{(1)} (\bar{q}_L^{(i)} H^{(1)}) q_{dR}^{(j)} + D_{ij}^{(2)} (\bar{q}_L^{(i)} H^{(2)}) q_{dR}^{(j)} + \\ &\quad + U_{ij}^{(1)} (\bar{q}_L^{(i)} \tilde{H}^{(1)}) q_{uR}^{(j)} + U_{ij}^{(2)} (\bar{q}_L^{(i)} \tilde{H}^{(2)}) q_{uR}^{(j)} + h.c. \} + \{\ \bar{\ell}\text{H}\ell\text{-terms}\ \}\ , \end{aligned} \tag{7}$$

where, however, the 3×3-Yukawa matrices $D^{(1)}$, $D^{(2)}$, $U^{(1)}$ and $U^{(2)}$ are now such that in the mass basis of quarks (fermions) they are all *simultaneously* diagonal. We note that "type I" and "type II" models (eqs.(1) and (2)) are special cases of this 2HD-SM(III). The corresponding charged-current part of the quarks can then be deduced

$$\mathcal{L}_{Y.}^{(III)cc} = H^{(+)}[-\overline{u_L} V D d_R + \overline{u_R} U^\dagger V d_L] + H^{(-)}[\text{ h.c. }]\ . \tag{8}$$

This expression is in the unitary gauge and in the physical bases (for fermions and the charged Higgses $H^{(\pm)}$). V is the CKM-matrix, $u^T = (u, c, t)$, $d^T = (d, s, b)$; U and D are specific linear combinations of the diagonal Yukawa matrices $U^{(j)}$ and $D^{(j)}$

$$U = -\sin\beta U^{(1)} + \cos\beta e^{-i\xi} U^{(2)}\ , \qquad D = -\sin\beta D^{(1)} + \cos\beta e^{+i\xi} D^{(2)}\ ,$$

where $\tan\beta$ is the ratio of the absolute values of the VEVs ($\tan\beta = v_2/v_1$), the mass basis is used, and ξ is the CP-violating phase between the VEVs. On the other hand, the (diagonal) mass matrices are obtained from the corresponding orthonormal linear combinations

$$\frac{\sqrt{2}}{v} M_u = \cos\beta U^{(1)} + \sin\beta e^{-i\xi} U^{(2)} \ , \qquad \frac{\sqrt{2}}{v} M_d = \cos\beta D^{(1)} + \sin\beta e^{+i\xi} D^{(2)} \ .$$

Therefore, assuming that no peculiar cancelations occur, U_{33} ($\sim (m_t^{phy}/v)$) is by far the most dominant Yukawa coupling in $\mathcal{L}_{Y.}^{(III)cc}$.

Since the $H^{(\pm)}$-exchanges influence the $\bar{B}_d^0$-B_d^0 and $\bar{K}^0$-K^0 mixing, the experimental values of $|m_{B^0} - m_{\bar{B}^0}|$, $|m_{K_L} - m_{K_S}|$ and of the CP-violating parameters ε_K and ε'_K would provide us with restrictions on the (dominant) U_{33} coupling strength of $H^{(\pm)}$ to quarks. The dominant graphs contributing to these mass differences are the W-W, H-W and H-H exchange box diagrams (cf. Fig. 4). The resulting effective 4-fermion couplings are

$$\begin{aligned} \mathcal{L}_{eff} (= \mathcal{L}_{eff}^{WW} + \mathcal{L}_{eff}^{HW} + \mathcal{L}_{eff}^{HH}) &\simeq \mathcal{A}^K [\overline{d(x)^a} \gamma^\mu (\frac{1-\gamma_5}{2}) s(x)^a]^2 \quad (\text{for } K^0 - \bar{K^0}) \\ &\simeq \mathcal{A}^B [\overline{b(x)^a} \gamma^\mu (\frac{1-\gamma_5}{2}) d(x)^a]^2 \quad (\text{for } B_d^0 - \bar{B_d^0}) \ , \end{aligned} \tag{9}$$

where: $\quad \mathcal{A} = \mathcal{A}_{WW} + \mathcal{A}_{HW} + \mathcal{A}_{HH} \ , \qquad \mathcal{A}_{HW} = \frac{G_F}{4\sqrt{2}\pi^2} \mathcal{J}(\frac{M_{H^-}}{M_W}, \frac{m_t}{M_W}) (\frac{m_t}{M_W})^2 \zeta_t^2 |U_{33}|^2 + \cdots \ ,$

$$\mathcal{A}_{HH} = -\frac{1}{64\pi^2 M_{H^-}^2} \mathcal{I}(\frac{m_t}{M_{H^-}}) \zeta_t^2 |U_{33}|^4 + \cdots \ , \qquad \mathcal{A}_{WW} = \frac{G_F^2 M_W^2}{4\pi^2} \{ \mathcal{E}(\frac{m_c}{M_W}) \zeta_c^2 + \cdots \} \ .$$

We denoted here

$$\zeta_t = V_{td}^* V_{ts} \ (\text{for } K^0 - \bar{K^0}) \ , \qquad \zeta_t = V_{td} V_{tb}^* \ (\text{for } \bar{B^0} - B^0) \ , \tag{10}$$

and analogously for ζ_c. $\mathcal{J}$, $\mathcal{I}$ and $\mathcal{E}$ are dimensionless and slowly varying functions ($\sim \mathcal{O}(1)$). Such effects, for the specific "type II" model, have been investigated by several authors [10]; here they are studied within the more general framework of the "type III" model. The relation of these amplitudes to the experimental inputs of K and B physics is

$$\begin{aligned} \sqrt{2} (\triangle M_{K_L^0 - K_S^0}) (|\varepsilon_K| + 0.05 \frac{\varepsilon'_K}{\varepsilon_K}) &\simeq -\frac{1}{2M_{K^0}} \mathrm{Im} \langle K^0 | \mathcal{L}_{eff}(x=0) | \bar{K^0} \rangle \\ &= -\frac{1}{2M_{K^0}} \mathcal{A}^K \mathrm{Im} \langle K^0 | [\overline{d^a} \gamma^\mu (\frac{1-\gamma_5}{2}) s^a]^2 | \bar{K^0} \rangle \ , \end{aligned} \tag{11}$$

$$M_{B^0} \triangle M_{B^0 - \bar{B}^0} \simeq | \langle \bar{B}_d^0 | \mathcal{L}_{eff}(x=0) | B_d^0 \rangle | = \mathcal{A}^B | \langle \overline{B_d^0} | [\bar{b^a} \gamma^\mu (\frac{1-\gamma_5}{2}) d^a]^2 | B_d^0 \rangle | \ , \tag{12}$$

where the normalization conventions are: $\langle P^0 | P^0 \rangle = Vol * 2M_{P^0}$ ($P^0 = K^0, B^0, \ldots$; Vol is the 3-dimensional volume of the space). While the box diagram amplitudes $\mathcal{A}^K$, $\mathcal{A}^B$ represent the dominant perturbative contributions in the above relations, the matrix elements $\mathrm{Im}\langle K^0 | \cdots | \bar{K^0} \rangle$ and $\langle \bar{B}_d^0 | \cdots | B_d^0 \rangle$ represent the hadronic (non-perturbative) effects. Experiments provide us with the following values for the $\triangle M$'s and the CP-violating parameters ε_K and ε'_K:

$$\begin{aligned} \sqrt{2} (\triangle M_{K_L^0 - K_S^0}) (|\varepsilon_K| + 0.05 \frac{\varepsilon'_K}{\varepsilon_K}) &\simeq 1.16 (1 \pm 0.03) 10^{-17} GeV \ , \\ \triangle M_{B^0 - \bar{B}^0} &\simeq (3.6 \pm 0.7) 10^{-13} GeV \ . \end{aligned} \tag{13}$$

For the hadronic matrix elements, we have the following (theoretical) uncertainties:

$$\langle K^0|[\bar{d^a}\gamma^\mu(\frac{1-\gamma_5}{2})s^a]^2|\bar{K^0}\rangle = \frac{2}{3}F_K^2 B_K M_{K^0}^2 \ , \qquad 0.6 \lesssim B_K \lesssim 0.9 \ .$$

$$\langle \bar{B_d^0}|[\bar{b^a}\gamma^\mu(\frac{1-\gamma_5}{2})d^a]^2|B_d^0\rangle = \frac{2}{3}F_B^2 B_B M_{B^0}^2 \ , \qquad 0.15GeV \lesssim F_B\sqrt{B_B} \lesssim 0.2GeV \ .$$

In addition, we have many uncertainties also in the CKM-matrix V (in $\mathcal{L}_{Y.}^{(III)cc}$). We chose Maiani's notation, fixing the real angles θ_{12}, θ_{13}, θ_{23} at their experimental average and allowing the CP-violating phase δ' in V to be free. Then, for a chosen m_t^{phy} and M_{H^-}, we obtain in the plane δ' vs. $|U_{33}|$ a stripe allowed by the hadronic uncertainties of the $\bar{B^0}$-B^0 mixing, and another stripe allowed by the hadronic uncertainties of the K^0-$\bar{K^0}$ mixing. For $m_t^{phy} = 175GeV$ and $M_{H^-} = 1TeV$, these stripes are depicted in Fig. 5, as well as their overlap (we use instead of $|U_{33}|$ the "normalized" parameter $Z = |U_{33}|/(\sqrt{2}m_t^{phy}/v)$). We included the leading (known) QCD corrections to the W-W exchange box diagrams, by making the following replacements in the amplitude $\mathcal{A}_{WW}^K$: $\zeta_c^2 \mapsto 0.81\zeta_c^2$, $\zeta_t^2 \mapsto 0.59\zeta_t^2$, $\zeta_t\zeta_c \mapsto 0.37\zeta_t\zeta_c$; and in $\mathcal{A}_{WW}^B$: $\zeta_t^2 \mapsto 0.8\zeta_t^2$. This changes the curves and the overlap quite substantially in the region $|U_{33}| \approx 0$. When decreasing M_{H^-} to $200GeV$, the features remain unchanged, but the upper bound for $|U_{33}|$ becomes smaller (Fig. 6). We note that it is the $\bar{B^0}$-B^0 mixing that provides us (for a given m_t^{phy} and M_{H^-}) with a rather restrictive upper bound on the dominant Yukawa coupling $|U_{33}|$. For heavier top quarks, this upper bound becomes even more restrictive (lower). Furthermore, in the special "type II" case, the horizontal axis in Figs. 5 and 6 should be interpreted as $\cot\beta$ ($= v_d/v_u$). In this case, Figs. 5 and 6 would give us lower bounds for $\tan\beta$.

Figs. 5 and 6: The regions in the δ' vs. Z ($= |U_{33}|/(2m_t^{phy}/v)$) parameter plane as allowed by the hadronic uncertainties for the case of $\bar{B_d^0}$-B_d^0 mixing (thick curves) and $\bar{K^0}$-K^0 mixing (thinner curves), for the case of the charged Higgs mass M_{H^-} = 1TeV (left) and M_{H^-} =0.2TeV (right). The overlap regions are also denoted. We took m_t^{phy} =175GeV, and took into account the leading QCD corrections for the W-W box exchange diagrams.

These rather preliminary calculations call for further improvements:

- the QCD corrections should also be included in the H-W and H-H exchange box diagrams (Fig. 4), thus improving Figs. 5 and 6 in the regime of large $|U_{33}|$;
- especially the experimentally uncertain θ_{13} angle of the CKM-matrix (in Maiani's notation) should be varied within the allowed range, thus giving us somewhat modified figures;

- similar calculations should be performed also for the "down-type" hadrons, e.g. for the D^0-$\bar{D}^0$ mixing, possibly giving us clues to restrictions on the dominant "down-type" coupling D_{33} of charged H^+ to b_R quark.

Conclusions

In general, standard models with two (or more) Higgs doublets (2HD-SMs) have a rich physics, especially concerning the CP violation [11]. However, there are also several other features that make these models attractive and very different from the minimal SM (MSM):

1. The 2HD-SM(II) has a clear and consistent convergence to flavor democracy as $E \uparrow \Lambda_{pole}$ (*unlike* the MSM and the closely related 2HD-SM(I)), thus giving us a signal of a possible new and (almost) flavor-blind strongly interacting physics at the energies beyond the Landau pole: $E \gtrsim \Lambda_{pole}$ ($\gtrsim 1TeV$).

2. For Λ_{pole} ($\sim E_{trans.}$) being smaller than E_{Planck}, it is possible to have $m_t^{phy} < 200GeV$ in the 2HD-SM(II) (*not* in the MSM and 2HD-SM(I)) $\Longrightarrow$ condensation scenarios leading to the 2HD-SM(II) (and to approximate flavor democracy at $E_{trans.} \sim E_{condens.} \sim \Lambda_{pole}$) may be promising.

3. The *imposed* mixed "lepton-quark" flavor democracy at $E \simeq \Lambda_{pole}$ in the 2HD-SM(II) leads to restricted values of the ratio of the VEVs ($0.53 < v_u/v_d < 2.1$, for $m_t^{phy} = (175 \pm 20)GeV$), and furthermore, would make the existence of the 4th generation of fermions very unlikely.

4. 2HD-SM(III) is the most general 2HD-SM with *no* Higgs-mediated flavor-changing neutral currents (FCNC) at the tree level, and is a promising framework for investigations; experimental FCNC (and CP violation) constraints, when confronted with 1-loop and higher loop effects, may severely restrict the parameter space of the model.

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Fig. 3: Bounds on m_t^{phy} as function of the VEV ratio v_u/v_d in the 2HD-SM(II); assumed: q-ℓ FD at Λ_{pole}, and see-saw.

Fig. 4: W-W, H-W and H-H exchange box diagrams for $\bar{P}^0$-P^0 mixing ($P^0 = B^0$, K^0).